

Stringy origin of Tevatron Wjj anomalyLuis A. Anchordoqui,¹ Haim Goldberg,² Xing Huang,¹
Dieter Lüst,^{3,4,5} and Tomasz R. Taylor²¹*Department of Physics,
University of Wisconsin-Milwaukee, Milwaukee, WI 53201, USA*²*Department of Physics,
Northeastern University, Boston, MA 02115, USA*³*Max-Planck-Institut für Physik
Werner-Heisenberg-Institut, 80805 München, Germany*⁴*Arnold Sommerfeld Center for Theoretical Physics
Ludwig-Maximilians-Universität München, 80333 München, Germany*⁵*Physics Division (TH), CERN 1211 Geneva 23, Switzerland
(Dated: April 2011)***Abstract**

The invariant mass distribution of dijets produced in association with W bosons, recently observed by the CDF Collaboration at Tevatron, reveals an excess in the dijet mass range $120 - 160 \text{ GeV}/c^2$, 3σ beyond Standard Model expectations. We show that such an excess is a generic feature of low mass string theory, due to the production and decay of a leptophobic Z' , a singlet partner of $SU(3)$ gluons coupled primarily to the $U(1)$ baryon number. In this framework, $U(1)$ and $SU(3)$ appear as subgroups of $U(3)$ associated with open strings ending on a stack of 3 D-branes. In addition, a minimal model contains two other stacks to accommodate the electro-weak $SU(2) \subset U(2)$ and the hypercharge $U(1)$. Of the three $U(1)$ gauge bosons, the two heavy Z' and Z'' receive masses through the Green-Schwarz mechanism. For a given Z' mass, the model is quite constrained. Fine tuning three of its free parameters is just sufficient to simultaneously ensure: a small $Z - Z'$ mixing in accord with the stringent LEP data on the Z mass; very small (less than 1%) branching ratio into leptons; and a large hierarchy between Z'' and Z' masses. The heavier neutral gauge boson Z'' is within the reach of LHC.

It appears that in the last year of the Tevatron's endeavors, it has pierced the Standard Model's resistant armor [1, 2]. The latest foray is an excess at $M_{jj} \simeq 140$ GeV in the dijet system invariant mass distribution of the associated production of a W boson with 2 jets (hereafter Wjj production) [2]. The CDF Collaboration fitted the excess to a Gaussian and estimated its production rate to be ~ 4 pb. This is roughly 300 times the Standard Model Higgs rate $\sigma(p\bar{p} \rightarrow WH) \times \text{BR}(H \rightarrow \bar{b}b)$. For a search window of 120 – 200 GeV, the excess significance above Standard Model background (including systematic uncertainties) is 3.2σ [2].

The CDF Wjj anomaly has been related to the technipion of a low mass technicolor [3], to resonant super-partner production in a supersymmetric model with R -parity violation [4], and to a leptophobic Z' gauge boson [5–8]. The suppressed coupling to leptons in the latter is required to evade the strong constraints of the Tevatron Z' searches in the dilepton mode [9]. All existing dijet-mass searches at the Tevatron are limited to $M_{jj} > 200$ GeV [10] and therefore cannot constrain the existence of a Z' with $M_{Z'} \simeq 140$ GeV. The strongest constraint on a light leptophobic Z' comes from the dijet search by the UA2 Collaboration, which has placed a 90% CL upper bound on $\sigma \times \text{BR}(Z' \rightarrow jj)$ in this energy range [11]. In this Letter we show that a Z' that can explain the Wjj excess and is in full agreement with existing limits on Z' coupling to quarks and leptons can materialize in the context of D-brane TeV-scale string compactifications.

At the time of its formulation and for years thereafter, superstring theory was regarded as a unifying framework for Planck-scale quantum gravity and TeV-scale Standard Model physics. Important advances were fueled by the realization of the vital role played by D-branes in connecting string theory to phenomenology. This has permitted the formulation of string theories with compositeness setting in at TeV scales and large extra dimensions [12].

TeV-scale superstring theory provides a brane-world description of the Standard Model, which is localized on membranes extending in $p + 3$ spatial dimensions, the so-called D-branes [13, 14]. Gauge interactions emerge as excitations of open strings with endpoints attached on the D-branes, whereas gravitational interactions are described by closed strings that can propagate in all nine spatial dimensions of string theory (these comprise flat parallel dimensions extended along the $(p + 3)$ -branes and transverse dimensions) [15]. The apparent weakness of gravity at energies below few TeV can then be understood as a consequence of the gravitational force “leaking” into the transverse compact dimensions of spacetime.

There are two peerless phenomenological consequences for TeV-scale D-brane string physics: the emergence of Regge recurrences at parton collision energies $\sqrt{\hat{s}} \sim M_s \equiv$ string scale, most distinctly manifest in the $\gamma + \text{jet}$ [16] and dijet [17] spectra resulting from their decay¹; and the presence of one or more additional $U(1)$ gauge symmetries, beyond the $U(1)_Y$ of the Standard Model. The latter follows from the property that the gauge group for open strings terminating on a stack of N identical D-branes is $U(N)$ rather than $SU(N)$ for $N > 2$. (For $N = 2$ the gauge group can be $Sp(1)$ rather than $U(2)$.)

To develop our program in the simplest way, we will work within the construct of a minimal model in which we consider scattering processes which take place on the (color) $U(3)$ stack of D-branes. In the bosonic sector, the open strings terminating on this stack contain, in addition to the $SU(3)$ octet of gluons g_μ^a , an extra $U(1)$ boson (C_μ , in the

¹ The recent search for string resonances in pp -collisions by the CMS collaboration [18] at the LHC now excludes a string scale below 2.5 TeV. From the string theory point of view, D-brane models with a TeV string scale can be obtained by the compactification on special Calabi-Yau spaces (Swiss cheese manifolds).

notation of [19] - this model was also discussed in [20]), most simply the manifestation of a gauged baryon number symmetry. The $U(1)_Y$ boson Y_μ , which gauges the usual electroweak hypercharge symmetry, is a linear combination of C_μ , the $U(1)$ boson B_μ terminating on a separate $U(1)$ brane, and perhaps a third additional $U(1)$ field X_μ sharing a $U(2)$ stack which is also a terminus for the $SU(2)_L$ electroweak gauge bosons W_μ^a [21]. Any vector boson Z'_μ , orthogonal to the hypercharge, must grow a mass in order to avoid long range forces between baryons other than gravity and Coulomb forces. The anomalous mass growth allows the survival of global baryon number conservation, preventing fast proton decay [22].

In the minimal $U(3) \times Sp(1) \times U(1)$ D-brane model, the hypercharge

$$Q_Y \equiv \frac{1}{6}Q_{U(3)} - \frac{1}{2}Q_{U(1)} \quad (1)$$

is anomaly free. However, the $Q_{U(3)}$ (gauged baryon number) is not anomaly free and we expect this anomaly to be canceled via a Green-Schwarz mechanism. There is an explicit mass term in the Lagrangian for the new gauge field $-\frac{1}{2}M'^2 Y'_\mu Y'^\mu$ whose scale comes from the compactification scheme. The scalar that gets eaten up to give the longitudinal polarization of the Y' is a closed string field and there is no extra Higgs particle [22]. Following [19] we take M' as a free parameter of the model and use precision electroweak data to determine its value. As usual, the $U(1)$ gauge interactions arise through the covariant derivative

$$\mathcal{D}_\mu = \partial_\mu - ig_1 B_\mu Q_{U(1)} - i\frac{g_3}{\sqrt{6}} C_\mu Q_{U(3)}, \quad (2)$$

where g_1 , g_2 , and g_3 are the gauge coupling constants. Introducing $S_P \equiv \sin \theta_P$ and $C_P \equiv \cos \theta_P$, the $U(1)$ fields can be projected into massless and massive directions

$$C_\mu = C_P Y'_\mu + S_P Y_\mu, \quad B_\mu = S_P Y'_\mu - C_P Y_\mu, \quad (3)$$

with

$$\tan \theta_P = \sqrt{\frac{2}{3}} \frac{g_1}{g_3}, \quad \text{and} \quad \frac{1}{g_{Y'}} = \frac{1}{6g_3^2} + \frac{1}{4g_1^2}. \quad (4)$$

Substituting (3) into (2) we obtain

$$g_{Y'} Q_{Y'} = \frac{g_3}{\sqrt{6}} C_P Q_{U(3)} + g_1 S_P Q_{U(1)}. \quad (5)$$

We note that a value for $g_{Y'}$ will emerge once a normalization for $Q_{Y'}$ is adopted. (The second relation in Eq. (4) depends on the choice of normalization for the hypercharge). For a Higgs ($Q_{U(3)} = 0$, $Q_{U(1)} = -1$, $Q_Y = -1/2$) with vacuum expectation value

$$\langle H \rangle = \begin{pmatrix} v \\ 0 \end{pmatrix}, \quad (6)$$

the kinetic term $(D_\mu H)^*(D_\mu H)$ gives a mass term

$$(v, 0) \begin{pmatrix} -\frac{1}{2}\sqrt{g_2^2 + g_Y^2} Z - g_1 S_P Y' & 0 \\ 0 & \frac{g_2^2 - g_Y^2}{2\sqrt{g_2^2 + g_Y^2}} Z - g_1 S_P Y' \end{pmatrix}^2 \begin{pmatrix} v \\ 0 \end{pmatrix} = (\overline{M}_Z Z + g_1 S_P v Y')^2, \quad (7)$$

TABLE I: Chiral fermion spectrum of the $U(3) \times U(2) \times U(1)$ D-brane model.

Name	Representation	$Q_{U(3)}$	$Q_{U(2)}$	$Q_{U(1)}$	Q_Y	$g_{Y'}Q_{Y'}$	$g_{Y''}Q_{Y''}$
U_i	$(\bar{3}, 1)$	2	0	0	$-\frac{4}{3}$	0.265	0.867
D_i	$(\bar{3}, 1)$	-1	0	1	$\frac{2}{3}$	-0.098	-0.444
L_i	$(1, 2)$	0	-1	1	-1	-0.004	-0.138
E_i	$(1, 1)$	0	2	0	2	0.078	0.255
Q_i	$(3, 2)$	1	1	0	$\frac{1}{3}$	0.172	0.561

where

$$\mathcal{D}_\mu = \partial_\mu - i \frac{1}{\sqrt{g_2^2 + g_Y^2}} Z_\mu (g_2^2 T^3 - g_Y^2 Y) - i g_{Y'} Y'_\mu Q_{Y'}, \quad (8)$$

with $T^3 = \sigma^3/2$ and $g_{Y'}Q_{Y'}$ given in Eq. (5). Equation (7) together with the mass term $\frac{1}{2}M'^2 Y'^2$ lead to a mass matrix

$$\frac{1}{2}(Z, Y') \begin{pmatrix} \overline{M}_Z^2 & \overline{M}_Z g_1 S_P v \\ \overline{M}_Z g_1 S_P v & g_1^2 S_P^2 v^2 + M'^2 \end{pmatrix} \begin{pmatrix} Z \\ Y' \end{pmatrix} = \frac{1}{2}(\overline{M}_Z Z + g_1 v S_P Y')^2 + \frac{1}{2}M'^2 Y'^2, \quad (9)$$

where $2\overline{M}_Z^2 = g_2^2 v^2 + g_Y^2 v^2$ is the usual tree level formula for the mass of the Z particle in the electroweak theory, before mixing [19]. When the theory undergoes electroweak symmetry breaking, because Y' couples to the Higgs, one gets additional mixing. However, to avoid conflict with precision measurements at LEP throughout this Letter we will enforce negligible $Z - Z'$ mixing and consider $M' \simeq M_{Z'}$ [23]. A comprehensive study of the M' parameter space has been carried out in [24], concluding that gauge bosons with $M_{Z'} < 700$ GeV are excluded by the Z -pole data from LEP.

In the $U(3) \times U(2) \times U(1)$ D-brane model the $Q_{U(1)}$, $Q_{U(2)}$, $Q_{U(3)}$ content of the hypercharge operator,

$$Q_Y = c_3 Q_{U(3)} + c_2 Q_{U(2)} + c_1 Q_{U(1)}, \quad (10)$$

is not uniquely determined by the anomaly cancellation requirement. In fact as seen in [21], there are 5 possibilities. This final choice does not depend on further symmetry considerations; in [21] it was fixed by requiring partial unification ($g_3 = g_2$) and acceptable value of $\sin^2 \theta_W$ at $M_s \sim 6 - 8$ TeV. We take $c_1 = 0$, $c_2 = 1$, and $c_3 = -2/3$ [25]. For this hypercharge embedding, conventional logarithmic running of coupling constants predicts a high string scale, $M_s \sim \mathcal{O}(10^{10})$ TeV [20]. However, it is possible that threshold corrections in the form of power law running can lower the string scale to the 5 – 10 TeV region [26]. In what follows we assume this to be the case. The chiral fermion spectrum² is summarized in Table I.

The covariant derivative is given by [27]

$$\mathcal{D}_\mu = \partial_\mu - i \frac{g_3}{\sqrt{6}} C_\mu Q_{U(3)} - i \frac{g_2}{2} X_\mu Q_{U(2)} - i g_1 B_\mu Q_{U(1)}. \quad (11)$$

² Note that we are considering D-brane quivers, where the right-handed up-quarks U_i originate from open strings ending on the $U(3)$ stack and its orientifold image. This leads to an antisymmetric matter representation with $Q_{U(3)} = 2$. Furthermore, the normalization of Q_Y for all matter fields in Table I is different from that in Eq.(1) by a factor of two.

The fields C_μ, X_μ, B_μ are related to Y_μ, Y'_μ and Y''_μ by a rotation matrix,

$$\mathcal{R} = \begin{pmatrix} C_\theta C_\psi & -C_\phi S_\psi + S_\phi S_\theta C_\psi & S_\phi S_\psi + C_\phi S_\theta C_\psi \\ C_\theta S_\psi & C_\phi C_\psi + S_\phi S_\theta S_\psi & -S_\phi C_\psi + C_\phi S_\theta S_\psi \\ -S_\theta & S_\phi C_\theta & C_\phi C_\theta \end{pmatrix}, \quad (12)$$

with Euler angles θ, ψ , and ϕ . Equation (11) can be rewritten in terms of Y_μ, Y'_μ , and Y''_μ as follows

$$\begin{aligned} \mathcal{D}_\mu = & \partial_\mu - iY_\mu \left(-S_\theta g_1 Q_{U(1)} + \frac{1}{2} C_\theta S_\psi g_2 Q_{U(2)} + \frac{1}{\sqrt{6}} C_\theta C_\psi g_3 Q_{U(3)} \right) \\ & - iY'_\mu \left[C_\theta S_\phi g_1 Q_{U(1)} + \frac{1}{2} (C_\phi C_\psi + S_\theta S_\phi S_\psi) g_2 Q_{U(2)} + \frac{1}{\sqrt{6}} (C_\psi S_\theta S_\phi - C_\phi S_\psi) g_3 Q_{U(3)} \right] \\ & - iY''_\mu \left[C_\theta C_\phi g_1 Q_{U(1)} + \frac{1}{2} (-C_\psi S_\phi + C_\phi S_\theta S_\psi) g_2 Q_{U(2)} + \frac{1}{\sqrt{6}} (C_\phi C_\psi S_\theta + S_\phi S_\psi) g_3 Q_{U(3)} \right]. \end{aligned} \quad (13)$$

Now, by demanding that Y_μ has the hypercharge Q_Y given in Eq. (10) we fix the first column of the rotation matrix \mathcal{R}

$$\begin{pmatrix} C_\mu \\ X_\mu \\ B_\mu \end{pmatrix} = \begin{pmatrix} Y_\mu \sqrt{6} c_3 g_Y / g_3 & \dots \\ Y_\mu 2c_2 g_Y / g_2 & \dots \\ Y_\mu c_1 g_Y / g_1 & \dots \end{pmatrix}, \quad (14)$$

and we determine the value of the two associated Euler angles

$$\theta = \arcsin[c_1 g_Y / g_1] = 0 \quad (15)$$

and

$$\psi = \arcsin[2c_2 g_Y / (g_2 C_\theta)] = 1.99, \quad (16)$$

where we have taken $M_Z = 91.1876$, $g_2 = 0.6596$, $g_3 = 1.215$. The third Euler angle ϕ and the coupling g_1 are determined by requiring sufficient suppression ($\lesssim 1\%$) to leptons and compatibility with the 90%CL upper limit reported by the UA2 Collaboration on $\sigma(p\bar{p} \rightarrow Z') \times \text{BR}(Z' \rightarrow jj)$ at $\sqrt{s} = 630$ GeV. The decay width of $Z' \rightarrow f\bar{f}$ is given by [30]

$$\Gamma(Z' \rightarrow f\bar{f}) = \frac{G_F M_Z^2}{6\pi\sqrt{2}} N_c C(M_{Z'}^2) M_{Z'} \sqrt{1-4x} [v_f^2(1+2x) + a_f^2(1-4x)], \quad (17)$$

where G_F is the Fermi coupling constant, $C(M_{Z'}^2) = 1 + \alpha_s/\pi + 1.409(\alpha_s/\pi)^2 - 12.77(\alpha_s/\pi)^3$, $\alpha_s = \alpha_s(M_{Z'})$ is the strong coupling constant at the scale $M_{Z'}$, $x = m_f^2/M_{Z'}^2$, v_f and a_f are the vector and axial couplings, and $N_c = 3$ or 1 if f is a quark or a lepton, respectively. The parton-parton cross section in the narrow Z' width approximation is given by

$$\hat{\sigma}(q\bar{q} \rightarrow Z') = K \frac{2\pi}{3} \frac{G_F M_Z^2}{\sqrt{2}} [v_q^2(\phi, g_1) + a_q^2(\phi, g_1)] \delta(\hat{s} - m_{Z'}^2), \quad (18)$$

where the K -factor represents the enhancement from higher order QCD processes estimated to be $K \simeq 1.3$ [31]. After folding $\hat{\sigma}$ with the CTEQ6 parton distribution functions [32],

taking $M_{Z'} = 140$ GeV, the branching ratio of electrons to quarks is minimized within the $\phi - g_1$ parameter space, subject to saturation of the 90%CL upper limit [11],

$$\sigma(p\bar{p} \rightarrow Z') \times \text{BR}(Z' \rightarrow jj) \approx 250 \text{ pb}. \quad (19)$$

This occurs for $\phi = 1.87$ and $g_1 = 0.036$, corresponding to a suppression $\Gamma_{Z' \rightarrow e^+e^-}/\Gamma_{Z' \rightarrow q\bar{q}} \sim 0.5\%$. (This also corresponds to $v_u^2 + a_u^2 = 0.355$, and $v_d^2 + a_d^2 = 0.139$.) The UA2 data has a dijet mass resolution $\Delta M_{jj}/M_{jj} \sim 10\%$ [11]. Therefore, at 140 GeV the dijet mass resolution is about 15 GeV. This is much larger than the resonance width, which is calculated to be $\Gamma(Z' \rightarrow f\bar{f}) \simeq 2$ GeV. All the couplings of the Y' boson are now determined and contained in Eq. (14). Numerical values are given in Table I under the heading of $g_{Y'}Q_{Y'}$. The corresponding Wjj production rate at the Tevatron ($\sqrt{s} = 1.96$ TeV) mediated through t and u channel quark exchange is found to be ≈ 4 pb, in agreement with observation [2] and with the recent estimate of [8]. The rate for the associated production channels ZZ' , $\gamma Z'$, and $Z'Z'$ is down by factors of approximately 3, 5, and 9, respectively [8].

The second strong constraint on the model derives from the mixing of the Z and the Y' through their coupling to the two Higgs doublets H and H' . The criteria we adopt here to define the Higgs charges is to make the Yukawa couplings ($H\bar{u}q$ and $H'\bar{d}q$) invariant under all three $U(1)$'s. This leads to $Q_{U(3)} = 3$, $Q_{U(2)} = 1$, $Q_{U(1)} = 0$, $Q_Y = -1$ and $Q_{U(3)} = 0$, $Q_{U(2)} = Q_{U(1)} = 1$, $Q_Y = 1$, for H and H' respectively.³ Here, $\langle H \rangle = \begin{pmatrix} v_u \\ 0 \end{pmatrix}$, $\langle H' \rangle = \begin{pmatrix} 0 \\ v_d \end{pmatrix}$, $v = \sqrt{v_u^2 + v_d^2} = 172$ GeV, and $\tan \beta \equiv v_u/v_d$ [29]. To account for Y'' we introduced a second term in (8), $\mathcal{D}_\mu = \partial_\mu \dots - ig_{Y'} Y'_\mu Q_{Y'} - ig_{Y''} Y''_\mu Q_{Y''}$, which is convenient to write as

$$-i \frac{x_H}{v_u} \overline{M}_Z Y'_\mu - i \frac{y_H}{v_u} \overline{M}_Z Y''_\mu + H \rightarrow H', \quad (20)$$

where for the two Higgs doublets

$$x_H = -0.252C_\phi + 1.886g_1S_\phi, \quad x_{H'} = 2.817C_\phi \quad (21)$$

and

$$y_H = 1.886g_1C_\phi + 0.252S_\phi, \quad y_{H'} = -2.817S_\phi. \quad (22)$$

The Higgs field kinetic term together with the Green-Schwarz mass terms ($-\frac{1}{2}M'^2 Y'_\mu Y'^\mu - \frac{1}{2}M''^2 Y''_\mu Y''^\mu$) yield the following mass square matrix

$$\begin{pmatrix} \overline{M}_Z^2 & \overline{M}_Z^2(x_H C_\beta^2 + x_{H'} S_\beta^2) & \overline{M}_Z^2(y_H C_\beta^2 + y_{H'} S_\beta^2) \\ \overline{M}_Z^2(x_H C_\beta^2 + x_{H'} S_\beta^2) & \overline{M}_Z^2(C_\beta^2 x_H^2 + S_\beta^2 x_{H'}^2) + M'^2 & \overline{M}_Z^2(C_\beta^2 x_H y_H + S_\beta^2 x_{H'} y_{H'}) \\ \overline{M}_Z^2(y_H C_\beta^2 + y_{H'} S_\beta^2) & \overline{M}_Z^2(C_\beta^2 x_H y_H + S_\beta^2 x_{H'} y_{H'}) & \overline{M}_Z^2(y_H^2 C_\beta^2 + y_{H'}^2 S_\beta^2) + M''^2 \end{pmatrix},$$

³ The Higgs fields H with $Q_{U(3)} = 3$ cannot simply be realized by a single open string in the considered D-brane quiver. It has to be thought as the antisymmetric product of three fundamental representations of $U(3)$. Alternatively we could have chosen the Higgs field H as an open string, corresponding to the bifundamental representation with charges $Q_{U(3)} = 0$, $Q_{U(2)} = Q_{U(1)} = -1$, $Q_Y = -1$. This minimal supersymmetric Standard Model quiver is also consistent with the constraints from string tadpole cancelation. However, the up-quark Yukawa couplings are then forbidden in string perturbation theory and must be generated through D-instanton effects.

where $x_H = 0.139$, $x_{H'} = -0.824$, $y_H = 0.221$, and $y_{H'} = -2.694$. The free parameters are $\tan\beta$, $M_{Z'}$, and $M_{Z''}$ which will be fixed by requiring the shift of the Z mass to lie within 1 standard deviation of the experimental value and $M_{Z'} = 140 \pm 2$ GeV. We are also minimizing $M_{Z''}$ to ascertain whether it can be detected at existing colliders. This leads to $\tan\beta = 0.4$, $M_{Z'} \simeq M' \simeq 140$ GeV, and $M_{Z''} \simeq M'' \geq 3$ TeV.

We now explore (at the parton level) prospects for searches of Z'' signals at the Large Hadron Collider (LHC). All the couplings of the Y'' boson are given in Table I under the heading of $g_{Y''}Q_{Y''}$. Using these figures we determine $\Gamma_{Z'' \rightarrow e^+e^-}/\Gamma_{Z'' \rightarrow q\bar{q}} \sim 0.7\%$. We therefore consider the standard bump-hunting procedure for dijet searches. We calculate a signal-to-noise ratio, with the signal rate estimated in the invariant mass window $[M_{Z''} - 2\Gamma, M_{Z''} + 2\Gamma]$. The noise is defined as the square root of the number of QCD background events in the same dijet mass interval for the same integrated luminosity. As an illustration, we take $M_{Z''} = 3$ TeV, for which $\Gamma(Z'' \rightarrow f\bar{f}) = 493$ GeV. For 10 fb^{-1} of data collected at $\sqrt{s} = 14$ TeV, we obtain a signal-to-noise ratio of 15σ .

An obvious question is whether the existing data allow determination of the string mass scale. The anomalous mass contributions to $M_{Z'}$ and $M_{Z''}$ are proportional (with computable coefficients [28]) to $g_{Y'}M_s$ and $g_{Y''}M_s$, respectively. However, existing data can only determine the products $g_{Y'}Q_{Y'}$ and $g_{Y''}Q_{Y''}$, see Table I. Therefore, a separate measurement of the different quark flavor charges (e.g., by tagging on b 's and t 's in Z'' decays) is necessary to determine the absolute normalization of the couplings and predict the string mass scale.

To summarize, we have considered a low-mass string compactification in which the Standard Model gauge multiplets originate in open strings ending on 3 D-branes. For the non-abelian $SU(3)$ and $SU(2)$ groups the D-brane construct requires the existence of two additional $U(1)$ bosons coupled to baryon number and to the trace of the $SU(2)$ multiplets, respectively. One linear combination of the three $U(1)$ gauge bosons is identified as the hypercharge Y field, coupled to the anomaly free hypercharge current. The two remaining linear combinations (Y', Y'') of the three $U(1)$'s are coupled to anomalous currents, and grow masses in accord with the Green-Schwarz mechanism. After electroweak breaking, mixing with the third component of isospin results in the three observable gauge bosons, where with small mixing $Z' \simeq Y'$, $Z'' \simeq Y''$.

For a fixed $M_{Z'}$, the model contain several free parameters – a single mixing angle and a gauge coupling constant unconstrained by the data – which are chosen to suppress the branching of Z' decay into leptons and to accommodate the UA2 90%CL data on $p\bar{p} \rightarrow jjX$. The remaining two parameters – $\tan\beta$ and $M_{Z''}$ – serve to limit the mass shift (due to mixing) of the electroweak Z to conform with LEP observations. The heavier neutral gauge boson Z'' is within the reach of LHC.

In closing, we note that there are some aspects of the model which can lead to observable consequences even in the absence of a light resonant signal. (1) The chiral nature of the couplings in Table I implies substantial parity violation. Hence, for $M_{Z'} \gtrsim 400$ GeV, the parity violating couplings of the Z' to fermions can generate a $t\bar{t}$ forward-backward asymmetry in $p\bar{p}$ collisions. (2) It was noted in [6] that both the Wjj anomaly and the forward-backward asymmetry observed at the Tevatron can be simultaneously explained by a Z' of $M_{Z'} \simeq 140$ GeV with flavor-violating coupling $g_{utZ'} \sim 0.45$. In principle these two conditions can be accommodated in D-brane constructions by introducing two quark families originating from strings stretching between two stacks of D-branes, and one family looping with both ends of a string attached to the color stack [13, 25]. This can give different charges to u and t quarks.

Acknowledgements

L.A.A. H.G. and T.R.T. are supported by the U.S. National Science Foundation (NSF Grant PHY-0757598 and PHY-0757959). D.L. likes to thank the theory department of CERN for its hospitality. The work was partially supported by the Cluster of Excellence "Origin and Structure of the Universe", in Munich. Any opinions, findings, and conclusions or recommendations expressed in this material are those of the authors and do not necessarily reflect the views of the National Science Foundation.

-
- [1] V. M. Abazov *et al.* [D0 Collaboration], Phys. Rev. Lett. **100**, 142002 (2008) [arXiv:0712.0851]; T. Aaltonen *et al.* [CDF Collaboration], Phys. Rev. Lett. **101**, 202001 (2008) [arXiv:0806.2472]; T. Aaltonen *et al.* [CDF Collaboration], arXiv:1101.0034.
 - [2] T. Aaltonen *et al.* [CDF Collaboration], arXiv:1104.0699.
 - [3] E. J. Eichten, K. Lane and A. Martin, arXiv:1104.0976.
 - [4] C. Kilic and S. Thomas, arXiv:1104.1002 [hep-ph].
 - [5] Y. Bai, B. A. Dobrescu, [arXiv:1012.5814 [hep-ph]].
 - [6] M. R. Buckley, D. Hooper, J. Kopp and E. Neil, arXiv:1103.6035.
 - [7] F. Yu, arXiv:1104.0243 [hep-ph].
 - [8] K. Cheung and J. Song, arXiv:1104.1375.
 - [9] D. E. Acosta *et al.* [CDF Collaboration], Phys. Rev. Lett. **95**, 131801 (2005) [arXiv:hep-ex/0506034]; T. Aaltonen *et al.* [CDF Collaboration], Phys. Rev. Lett. **99**, 171802 (2007) [arXiv:0707.2524]; T. Aaltonen *et al.* [CDF Collaboration], Phys. Rev. Lett. **102**, 091805 (2009) [arXiv:0811.0053].
 - [10] F. Abe *et al.* [CDF Collaboration], Phys. Rev. Lett. **74**, 3538 (1995) [arXiv:hep-ex/9501001]; F. Abe *et al.* [CDF Collaboration], Phys. Rev. D **55**, 5263 (1997) [arXiv:hep-ex/9702004]; B. Abbott *et al.* [D0 Collaboration], Phys. Rev. Lett. **82**, 2457 (1999) [arXiv:hep-ex/9807014]; T. Aaltonen *et al.* [CDF Collaboration], Phys. Rev. D **79**, 112002 (2009) [arXiv:0812.4036].
 - [11] J. Alitti *et al.* [UA2 Collaboration], Z. Phys. C **49**, 17 (1991); J. Alitti *et al.* [UA2 Collaboration], Nucl. Phys. B **400**, 3 (1993).
 - [12] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, Phys. Lett. B **436**, 257 (1998) [arXiv:hep-ph/9804398].
 - [13] R. Blumenhagen, B. Körs, D. Lust, T. Ott, Nucl. Phys. B **616**, 3 (2001) [hep-th/0107138].
 - [14] G. Honecker, T. Ott, Phys. Rev. D **70**, 126010 (2004) [hep-th/0404055]; F. Gmeiner, G. Honecker, JHEP **0807**, 052 (2008) [arXiv:0806.3039 [hep-th]].
 - [15] For a review see R. Blumenhagen, B. Körs, D. Lust, S. Stieberger, Phys. Rept. **445**, 1 (2007) [hep-th/0610327].
 - [16] L. A. Anchordoqui, H. Goldberg, S. Nawata and T. R. Taylor, Phys. Rev. Lett. **100**, 171603 (2008) [arXiv:0712.0386]; L. A. Anchordoqui, H. Goldberg, S. Nawata and T. R. Taylor, Phys. Rev. D **78**, 016005 (2008) [arXiv:0804.2013].
 - [17] L. A. Anchordoqui, H. Goldberg, D. Lust, S. Nawata, S. Stieberger and T. R. Taylor, Phys. Rev. Lett. **101**, 241803 (2008) [arXiv:0808.0497]; L. A. Anchordoqui, H. Goldberg, D. Lust, S. Nawata, S. Stieberger and T. R. Taylor, Nucl. Phys. B **821**, 181 (2009) [arXiv:0904.3547].
 - [18] V. Khachatryan *et al.* [CMS Collaboration], Phys. Rev. Lett. **105**, 211801 (2010) [arXiv:1010.0203 [hep-ex]].

- [19] D. Berenstein and S. Pinansky, Phys. Rev. D **75**, 095009 (2007) [arXiv:hep-th/0610104].
- [20] P. Anastasopoulos, T. P. T. Dijkstra, E. Kiritsis and A. N. Schellekens, Nucl. Phys. B **759**, 83 (2006) [arXiv:hep-th/0605226].
- [21] I. Antoniadis, E. Kiritsis and T. N. Tomaras, Phys. Lett. B **486**, 186 (2000) [arXiv:hep-ph/0004214].
- [22] D. M. Ghilencea, L. E. Ibanez, N. Irges and F. Quevedo, JHEP **0208**, 016 (2002) [arXiv:hep-ph/0205083].
- [23] Y. Umeda, G. C. Cho and K. Hagiwara, Phys. Rev. D **58**, 115008 (1998) [arXiv:hep-ph/9805447].
- [24] D. Berenstein, R. Martinez, F. Ochoa and S. Pinansky, Phys. Rev. D **79**, 095005 (2009) [arXiv:0807.1126].
- [25] D. Lüst, S. Stieberger and T. R. Taylor, Nucl. Phys. B **808**, 1 (2009) [arXiv:0807.3333].
- [26] K. R. Dienes, E. Dudas and T. Gherghetta, Phys. Lett. B **436**, 55 (1998) [arXiv:hep-ph/9803466]; K. R. Dienes, E. Dudas and T. Gherghetta, Nucl. Phys. B **537**, 47 (1999) [arXiv:hep-ph/9806292].
- [27] L. A. Anchordoqui, W. Z. Feng, H. Goldberg, X. Huang and T. R. Taylor, arXiv:1012.3466.
- [28] I. Antoniadis, E. Kiritsis and J. Rizos, Nucl. Phys. B **637**, 92 (2002) [arXiv:hep-th/0204153].
- [29] I. Antoniadis, E. Kiritsis, J. Rizos and T. N. Tomaras, Nucl. Phys. B **660**, 81 (2003) [arXiv:hep-th/0210263].
- [30] V. D. Barger, K. M. Cheung and P. Langacker, Phys. Lett. B **381**, 226 (1996) [arXiv:hep-ph/9604298].
- [31] V. Barger and R. J. N. Phillips, *Collider Physics* (Addison-Wesley, 1987).
- [32] J. Pumplin, D. R. Stump, J. Huston, H. L. Lai, P. Nadolsky and W. K. Tung, JHEP **0207**, 012 (2002) [arXiv:hep-ph/0201195].